A unified, flavor symmetric explanation for the $t\bar{t}$ asymmetry and Wjj excess at CDF

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We present a simple, perturbative, and renormalizable model with a flavor symmetry which can explain both the $t\bar{t}$ forward-backward asymmetry and the bump feature present in the dijet mass distribution of the $W+jj$ sample in the range 120–160 GeV that was recently reported by the CDF collaboration. The flavor symmetry not only ensures the flavor/CP safety of the model, but also relates the two anomalies unambiguously. It predicts a comparable forward-backward asymmetry in $c\bar{c}$. The forward-backward asymmetry in $b\bar{b}$ is, however, small. A bump in the dijet mass distribution in $Z+jj$ sample is also predicted but with a suppressed cross-section.

I. INTRODUCTION

Recently, the CDF collaboration has reported two interesting anomalies — a large $t\bar{t}$ forward-backward (FB) asymmetry [1, 2], and a 3.2σ excess in the 120–160 GeV range in the dijet mass distribution of the $W+jj$ sample [3] (see, however, Ref. [4] for a DØ analysis). The recent report of the FB asymmetry also confirms the trend suggested by the earlier measurements by CDF [5, 6] and D0 [7, 8].

It is a straightforward exercise to fit these two anomalies by introducing new particles with appropriately chosen masses and couplings. However, the nature of these anomalies suggests that the new physics should couple to standard-model (SM) quarks at tree level with an $\mathcal{O}(0.1)$–$\mathcal{O}(1)$ coupling along with a nontrivial quark flavor structure. Such a new physics typically faces strong constraints from precision flavor and CP constraints, unless the model is equipped with a flavor symmetry.

In this paper, we present a weakly-coupled, renormalizable field theory with a flavor symmetry to explain both anomalies. (For attempts to generate just the FB asymmetry preserving the full flavor symmetries, see e.g. Refs. [9–12].) We introduce just one multiplet of scalars with a single coupling to SM quarks dictated by the flavor symmetry. The flavor symmetry ensures that the only source of flavor and CP violations is $V_{\rm CKM}$. It also relates the sizes of the two anomalies in a definite manner, and entails additional predictions.

The rest of the paper is organized as follows. In Sec. II, we define our model with an emphasis on the flavor symmetry structure, which keeps flavor/CP violations under control without tuning or an ad hoc choice of couplings. In Sec. III, we go through various potential constraints on the model for the values of parameters necessary for obtaining the $t\bar{t}$ FB asymmetry and Wjj bump. Sec. IV shows our estimation of the asymmetry, while Sec. V shows the details of the bump feature in the dijet mass spectrum as predicted in our model. Our concluding reflections and a brief discussion of various implications of the model are included in Sec. VI.

II. THE MODEL

The flavor symmetry we propose is a subgroup of the $\mathrm{U}(3)^3$ quark flavor symmetry of the SM:

$$\left(\prod_{i=1}^{3} \mathrm{U}(1)_{q_{\mathrm{L}i}} \times \mathrm{U}(1)_{u_{\mathrm{R}i}}\right) \times \mathrm{U}(3)_d \times \mathbb{Z}_3\,, \qquad (1)$$

where $q_{\mathrm{L}i}$ and $u_{\mathrm{R}i}$ have charge $+1$ under $\mathrm{U}(1)_{q_{\mathrm{L}i}}$ and $\mathrm{U}(1)_{u_{\mathrm{R}i}}$, respectively, while d_{R} is a $\mathbf{3}$ of $\mathrm{U}(3)_d$. $\mathbb{Z}_3$ cyclically permutes the flavor indices of $q_{\mathrm{L}i}$ and $u_{\mathrm{R}i}$ ($i = 1, 2, 3$), but not of $d_{\mathrm{R}i}$. The lepton sector of our model is identical to that of the SM, and will not be discussed in this paper.

In the SM, one can always go to a basis where Y_u is diagonal. In the limit of neglecting both Y_u and Y_d, the SM possesses the flavor symmetry (1). Turning on the diagonal (but non-degenerate) Y_u breaks the symmetry (1) to its subgroup $\mathrm{U}(1)_{B_1} \times \mathrm{U}(1)_{B_2} \times \mathrm{U}(1)_{B_3} \times \mathrm{U}(3)_d$, where $q_{\mathrm{L}i}$ and $u_{\mathrm{R}i}$ have charge $+1$ under $\mathrm{U}(1)_{B_i}$. This subgroup still forbids all flavor violations. Turning on Y_d then breaks $\mathrm{U}(1)_{B_1} \times \mathrm{U}(1)_{B_2} \times \mathrm{U}(1)_{B_3} \times \mathrm{U}(3)_d$ down to the baryon number $\mathrm{U}(1)_B$, thus introducing flavor mixing. However, since Y_d breaks (and only Y_d breaks) $\mathrm{U}(3)_d$, we can always bring Y_d into the form $Y_d = V_{\rm CKM}\,\mathrm{diag}(y_d, y_s, y_b)$, ensuring that $V_{\rm CKM}$ is the only source of flavor violation. Also, note that $\mathrm{diag}(y_u, y_c, y_t)$ and $\mathrm{diag}(y_d, y_s, y_b)$ can both be taken to be positive definite, rendering $V_{\rm CKM}$ the only source of CP violation.[1]

Our fundamental assumption is that this symmetry breaking pattern persists for new physics beyond the SM as well. In other words, new physics should fully respect the flavor symmetry (1) in the limit $Y_u, Y_d \to 0$, and so Y_u and Y_d remain the only spurions breaking the symmetry (1). This can be thought of as a variant of minimal flavor violation (MFV) [13–17], and, in particular, the breaking pattern $\mathrm{U}(1)_{B_1} \times \mathrm{U}(1)_{B_2} \times \mathrm{U}(1)_{B_3} \times \mathrm{U}(3)_d \longrightarrow U(1)_B$ by Y_d has been studied in the context

[1] We neglect the QCD vacuum angle. It is straightforward to add an axion to our model to solve the strong CP problem.

of the supersymmetric SM [18, 19].[2] We introduce a $\mathbb{Z}_3$ triplet of complex scalar fields, $\Phi = (\Phi_1, \Phi_2, \Phi_3)$, where the gauge quantum number of Φ is $(\mathbf{1}, \mathbf{2})_{-1/2}$ under $(\mathrm{SU}(3)_\mathrm{C}, \mathrm{SU}(2)_\mathrm{L})_{\mathrm{U}(1)_\mathrm{Y}}$ representation. Φ_i $(i = 1, 2, 3)$ are singlets under $\mathrm{U}(3)_d$, but they are charged under $\mathrm{U}(1)_{q_{\mathrm{L}1}} \times \mathrm{U}(1)_{q_{\mathrm{L}2}} \times \mathrm{U}(1)_{q_{\mathrm{L}3}}$ as

$$\Phi_1 \sim (0,0,1)\,,\ \Phi_2 \sim (1,0,0)\,,\ \Phi_3 \sim (0,1,0)\,. \qquad (2)$$

while under $\mathrm{U}(1)_{u_{\mathrm{R}1}} \times \mathrm{U}(1)_{u_{\mathrm{R}2}} \times \mathrm{U}(1)_{u_{\mathrm{R}3}}$ as

$$\Phi_1 \sim (0,-1,0)\,,\ \Phi_2 \sim (0,0,-1)\,,\ \Phi_3 \sim (-1,0,0)\,, \qquad (3)$$

Note that these charge assignments respect $\mathbb{Z}_3$. The tree-level Lagrangian reads:

$$\begin{aligned}\mathcal{L}_\mathrm{tree} = {} & \mathcal{L}_\mathrm{SM} + (D_\mu\Phi)^\dagger(D^\mu\Phi) - m_\Phi^2\Phi^\dagger\Phi \\ & - \lambda\,(\overline{q}_{\mathrm{L}1}\Phi_2 u_{\mathrm{R}3} + \overline{q}_{\mathrm{L}2}\Phi_3 u_{\mathrm{R}1} + \overline{q}_{\mathrm{L}3}\Phi_1 u_{\mathrm{R}2} + \mathrm{c.c.}) \\ & - \xi\,(H^\dagger\sigma^a H)(\Phi^\dagger\sigma^a\Phi) - \xi'(H^\dagger H)(\Phi^\dagger\Phi) \\ & - \frac{\zeta}{4}(\Phi^\dagger\Phi)^2\,.\end{aligned} \qquad (4)$$

The new physics is completely invariant under the flavor symmetry (1). Therefore, the flavor symmetry (1) is broken only by Y_u and Y_d, just as in the SM. Also, we have chosen λ to be real without loss of generality, by redefining the phase of Φ. No new CP phase has, therefore, been introduced in this model.

The tree-level Lagrangian (4) has three phenomenologically relevant parameters: m_Φ, λ, and ξ. $\mathbb{Z}_3$ dictates that all three components of Φ have the equal mass m_Φ, and that they all couple to the SM quarks with the same strength λ and likewise to the Higgs via ξ and ξ'. The interesting role of ξ is that it splits the masses of the neutral (Φ^0) and charged (Φ^-) components of Φ as:

$$m_{\Phi^0}^2 = m_{\Phi,\mathrm{eff}}^2 - \xi v^2\,, \quad m_{\Phi^-}^2 = m_{\Phi,\mathrm{eff}}^2 + \xi v^2\,, \qquad (5)$$

where $v = 174$ GeV and $m_{\Phi,\mathrm{eff}}^2 \equiv m_\Phi^2 + \xi' v^2$. We choose $m_{\Phi^0} = 160$ GeV and $m_{\Phi^-} = 220$ GeV, which corresponds to $m_{\Phi,\mathrm{eff}} = 192$ GeV and $\xi = 0.38$. We take λ to be 1.4. This might appear too large to keep λ perturbative up to very high scale. Fortunately, the one-loop RG equation for λ is similar to that of the top Yukawa coupling in the SM and there is a quasi-fixed point near $\lambda \approx 1.4$.

At loop level, counter-terms $\delta\mathcal{L}$ must be added to the Lagrangian (4) for renormalization. We assume that all terms required to renormalize the theory are present in $\delta\mathcal{L}$, at the minimal level required to avoid fine tuning.[3] This assumption is technically natural, and may be justified by assuming that our Lagrangian arises from a more fundamental theory in which Y_u and Y_d are the only parameters breaking the flavor symmetry (1). As an example of terms in $\delta\mathcal{L}$, renormalization requires the counterterms $\overline{q}_{\mathrm{L}1} Y_{u1} \Phi_3^\dagger (Y_d d_\mathrm{R})_2$ + (cyclic permutations) with a common coefficient $\sim \lambda/(16\pi^2)\log\Lambda$. So we include these operators in $\delta\mathcal{L}$ with a single coefficient $\sim \lambda/(16\pi^2)$. This also exemplifies the general principle that all generated operators respect the flavor symmetry (1), broken only by Y_u and Y_d. Moreover, since no operators are generated with an independent phase, renormalization does not require introducing new phases. The property that V_CKM is the only source of flavor and CP violations, therefore, remains intact at the quantum level.

III. CONSTRAINTS

The property that V_CKM is the only source of flavor/CP violations, and the fact that the mass scale for Φ^0 and Φ^- is similar to or larger than the top quark mass, imply that flavor/CP violations involving Φ is at most comparable to those in the SM. For example, consider bounds from D^0-$\overline{D}^0$ mixing, that is, 4-fermion operators with two c and two $\bar{u}$ fields that arise upon integrating Φ out. Recall that, in our flavor symmetry breaking pattern, the only flavor non-diagonal spurion is $Y_d = V_\mathrm{CKM}\,\mathrm{diag}(y_d, y_s, y_b)$. The flavor symmetry (1) thus dictates that the simplest combination of spurions that can change c to u must involve the combination $(Y_d Y_d^\dagger)_{12} = (V_\mathrm{CKM}\,\mathrm{diag}(y_d^2, y_s^2, y_b^2) V_\mathrm{CKM}^\dagger)_{12} \sim \mathcal{O}(10^{-6})$. Since this combination converts c_L to s_L, the coefficient of $(\bar{u}_\mathrm{L} c_\mathrm{L})^2$ is on the order of $\left[(Y_d Y_d^\dagger)_{12}/m_\Phi\right]^2 \sim (10^5\ \mathrm{TeV})^{-2}$ multiplied by certain powers of the couplings λ and g and the loop factor $1/(16\pi^2)$. This is safely much smaller than the experimental bound $\sim (10^3\ \mathrm{TeV})^{-2}$ [20]. For the $(\bar{u}_\mathrm{L} c_\mathrm{R})(\bar{u}_\mathrm{R} c_\mathrm{L})$ operator, the spurions $(Y_u)_{22} \sim 10^{-2}$ and $(Y_u)_{11} \sim 10^{-5}$ have to be further inserted to convert c_L to c_R and u_L to u_R, respectively, rendering the coefficient of the operator way below the bound $\sim (10^4\,\mathrm{TeV})^{-2}$ [20]. Therefore, D^0-$\overline{D}^0$ mixing is not an issue at all in our model, thanks to the flavor symmetry.

The most stringent flavor bounds on our model arise from the 4-fermion operators that are generated via the tree-level exchange of Φ. In the gauge basis, these are

$$(\bar{q}_{\mathrm{L}2} u_\mathrm{R})(\bar{u}_\mathrm{R} q_{\mathrm{L}2})\,,\ (\bar{q}_{\mathrm{L}3} c_\mathrm{R})(\bar{c}_\mathrm{R} q_{\mathrm{L}3})\,,\ (\bar{q}_{\mathrm{L}1} t_\mathrm{R})(\bar{t}_\mathrm{R} q_{\mathrm{L}1})\,. \quad (6)$$

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[2] A crucial difference between MFV and our flavor symmetry breaking pattern is in the spurion structure in the up-quark sector. In MFV there are 9 complex spurions for the up-quark sector, that is, the 3×3 matrix Y_u transforming as $(\mathbf{3}, \bar{\mathbf{3}})$ under $\mathrm{SU}(3)_{q_\mathrm{L}} \times \mathrm{SU}(3)_{u_\mathrm{R}}$, while in our case there are only three real spurions, $(Y_u)_{ii}$ $(i = 1, 2, 3)$ carrying the charge $(1, -1)$ for $\mathrm{U}(1)_{q_{\mathrm{L}i}} \times \mathrm{U}(1)_{u_{\mathrm{R}i}}$. By assumption Y_d is the only spurion which breaks CP or the $\mathrm{U}(1)_{B_{1,2,3}}$ quantum numbers. Note that the up-type quarks are in the mass basis at the outset and that, unlike in MFV, unitary rotations done on q_L or u_R are not approximate symmetries of the Lagrangian.

[3] We have the usual "hierarchy problem" for m_H and m_Φ as well as other dimension-2 operators in $\delta\mathcal{L}$. In this paper, we simply tune m_H and m_Φ and use dimensional regularization with (modified) minimal subtraction to obtain the natural sizes of dimension-2 operators in $\delta\mathcal{L}$, but it is straightforward to supersymmetrize the model to justify this.

The first and second operators can contribute to hadronic b decays. In the mass basis, they contain

$$V_{cb}^* V_{ci} (\bar{b}_{\rm L} u_{\rm R})(\bar{u}_{\rm R} d_{{\rm L}i}) + V_{tb}^* V_{ti} (\bar{b}_{\rm L} c_{\rm R})(\bar{c}_{\rm R} d_{{\rm L}i})\,. \quad (7)$$

This comes from the tree-level exchange of Φ^-, so its coefficient is $\lambda^2/m_{\Phi^-}^2$. The first operator contributes to the charmless process $b \to s\bar{u}u$. The particle data book [21] specifies the total inclusive branching fraction of B mesons into charmed modes to be $(95 \pm 5)\%$. For $m_{\Phi^-} = 220$ GeV and $\lambda = 1.4$, the leading order spectator decay model gives a branching fraction for the $b \to s\bar{u}u$ mode of 15%, which is within the 2σ margin of error. CP constraints do not pose a problem for the new contribution to decays. The CP phase in the new contribution to $b \to c\bar{c}s$ is almost the same as the phase of the standard model contribution. The phase of the new contribution to the $b \to s\bar{u}u$ mode is the same as the gluonic penguin contribution, which so far, is consistent with experiments. Of greater concern is the nonstandard increase in the hadronic width. However the heavy quark expansion, which is needed for a theoretical computation of the hadronic width [22–24], may have significant uncertainty for this computation [25, 26]. Should further experimental tests of B decay modes exclude such a large new contribution to hadronic B decays then the Φ^- mass would have to be increased.

The correction to the electroweak parameter αT is given by

$$\alpha T = \frac{3}{32\pi^2 v^2}\left[m_{\Phi^0}^2 + m_{\Phi^-}^2 - \frac{2 m_{\Phi^-}^2 m_{\Phi^0}^2}{m_{\Phi^-}^2 - m_{\Phi^0}^2} \log \frac{m_{\Phi^-}^2}{m_{\Phi^0}^2} \right] \quad (8)$$

where α is to be evaluated at the weak scale. For $m_{\Phi^0} =$ 160 GeV and $m_{\Phi^-} = 220$ GeV, we get $\alpha T = 1.5 \times 10^{-3}$. From the particle data book [21], for a Higgs mass of 117(300) GeV, the T parameter is constrained to be $T = 0.07(0.16) \pm 0.08$. Thus, our model is consistent with precision electroweak constraints without any tuning. However, the T parameter contribution gives an upper bound on the mass of the Φ^-.

Φ couplings to leptons are highly loop-suppressed; so precision lepton measurements (e.g. the muon $g-2$) do not place constraints on our model. Since all components of Φ are unstable even in the collider time scale, there are no cosmological constraints. Hence, we concentrate on the collider physics constraints below.

- The total $t\bar{t}$ production cross-section is not changed significantly. Also note that the precise value of the theoretical prediction is still open to debate. The NLO+NNL calculations quote $\sim 10\%$ uncertainty [27–30] and the resummations of threshold logs result in a smaller value [31, 32]. When compared to the SM leading order result, we find that $d\sigma_{t\bar{t}}/dM_{t\bar{t}}$, where $M_{t\bar{t}}$ is the invariant mass of the $t\bar{t}$ pair, in our model slightly decreases near the threshold but slightly increases at higher values of $M_{t\bar{t}}$. (See Sec. IV for details).

- Single top production via Φ is suppressed. The leading single top production comes from $b\bar{c} \to \Phi_1^- \to W^- \Phi_1^0 \to W^- t\bar{c}$, which is highly suppressed by the smallness of the b and c parton distribution functions (PDFs). It is also suppressed by the 3-body phase space if the W^- is on-shell and the Φ^- is off-shell. If the Φ^- is on-shell, then the W^- must be off-shell, becoming a 4-body process. Either case, it is clearly much smaller than the SM counterpart $u\bar{d} \to W^+ \to t\bar{b}$, which is only suppressed by the offshellness of the W^+.

- The CDF collaboration has searched for events with same-sign lepton pairs and at least one b jet, and found 3 such events in 2 fb^{-1} of data [33], where they expect ~ 2 events from background. Di-top (tt) production can give such final states and is, thus, severely constrained. In our model, however, di-top production is extremely suppressed since in the limit of neglecting Y_d, baryon numbers for the three up-type quarks are separately conserved as in Eq. (1). Therefore, this is not a constraint for us.

- The top-quark width is not modified significantly. For $m_{\Phi^0} = 160$ GeV and $\lambda = 1.4$, the total top width is 1.6 GeV, which is well within the experimental limit [34, 35].

- There is a sizeable dijet production in our model via Φ_3 exchange in the s- or t-channel. For masses as low as 160 GeV or 220 GeV, Tevatron has large SM dijet backgrounds due to the gluon PDF's increasing rapidly at low parton x. Thus, there are no constraints from Tevatron [37]. The strongest bound comes from the CERN SPS collider ($p\bar{p}$ at $\sqrt{s} = 630$ GeV). The UA2 collaboration at the SPS has placed 90% C.L. bounds ≈ 100 pb on a dijet resonance at 160 GeV, and ≈ 10 pb at 220 GeV [36]. Among our dijet channels, those from t-channel Φ_3 exchange give a smooth dijet mass distribution on top of the smooth huge SM background. So they could not have been picked up by the UA2 search. For the s-channel contributions from Φ_3^0 and Φ_3^-, we find the cross-sections to be $\simeq 56$ pb for Φ_3^0 and $\simeq 15$ pb for Φ_3^-, with $\lambda = 1.4$. The latter may appear to be in conflict with the UA2 90% C.L. bound. However, note that the UA2 bounds are for a *narrow* resonance such as W' which can distinguish itself from the smooth SM dijet background. Our Φ_3^- resonance, on the other hand, is not narrow — its width is quite large, ≈ 26 GeV for $m_{\Phi^-} = 220$ GeV and $\lambda = 1.4$, which is expected to be smoothed out even further once parton showering and detector effects are taken into account. Thus, the search optimized for a very narrow resonance has a reduced sensitivity to our Φ_3^-. Furthermore, note that UA2 performed their analysis in the early days of QCD jet studies. Their answer depends crucially on the quality of the Monte Carlo

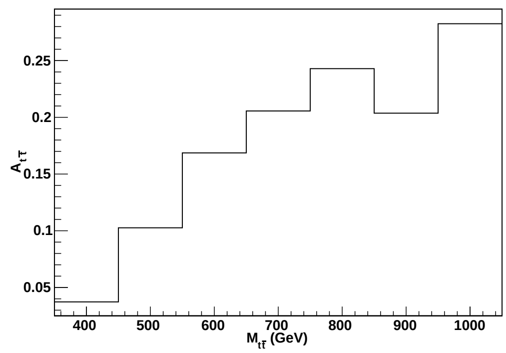

FIG. 1: $A_{t\bar{t}}$ as a function of the invariant mass of the $t\bar{t}$ pair, $M_{t\bar{t}}$. We calculated $A_{t\bar{t}}$ after demanding both the tops to be within $\eta = \pm 2.0$.

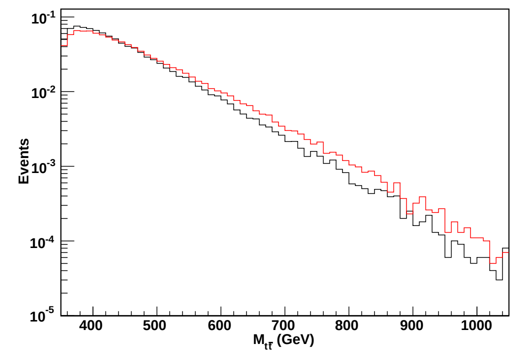

FIG. 2: Comparison of the area normalized distribution of $\sigma_{t\bar{t}}$ in the SM (in black) and in our model (in red). In both cases the cross-section is calculated at the leading order using `Madgraph v4.4.48`.

and the detector simulation which are primitive by today's standard. They also use events with two exclusive jets, where jets were constructed using an infrared unsafe jet algorithm [38]. We believe, therefore, that there are considerable uncertainties associated with their bounds, and that it is fair to regard the UA2 90% C.L. bound as an order-of-magnitude limit.

IV. THE $t\bar{t}$ FORWARD-BACKWARD ASYMMETRY

Neglecting the CKM mixings and non-valence partons in the incoming p and $\bar{p}$, the relevant interactions for this are:

$$\mathcal{L}_{\rm int} = -\lambda(\bar{u}_{\rm L}\Phi_2^0 t_{\rm R} + \bar{d}_{\rm L}\Phi_2^- t_{\rm R} + {\rm c.c.})\,. \qquad (9)$$

The $t\bar{t}$ forward-backward asymmetry arises at the Tevatron from the processes $u\bar{u} \to t\bar{t}$ with a t-channel Φ_2^0 exchange, and $d\bar{d} \to t\bar{t}$ with a t-channel Φ_2^- exchange.

A dedicated simulation including parton showering and detector effects is beyond the scope of this paper, partly due to the uncertainties in the SM prediction and the experimental measurement of the asymmetry. We simply perform our analysis at the parton level, and show that the asymmetry is generated with the right sign and is of the same order in magnitude.

We define the asymmetry as

$$A_{t\bar{t}} = \frac{N(\Delta y > 0) - N(\Delta y < 0)}{N(\Delta y > 0) + N(\Delta y < 0)}\,, \qquad (10)$$

where Δy is the rapidity difference between the t and $\bar{t}$. Our $t\bar{t}$ sample was generated using `Madgraph v4.4.48` [39]. We have imposed the following cuts on the $t\bar{t}$ pairs: $|\eta_t|, |\eta_{\bar{t}}| < 2.0$, and $M_{t\bar{t}} > 450$ GeV. We find the asymmetry to be $A_{t\bar{t}} \simeq 0.13$. Note that the asymmetry in our model does not depend linearly on $\cos\theta$ as it is assumed in Ref. [1] to extrapolate the asymmetry to the full 4π solid angle. This is because our asymmetry is generated by t-channel exchange of particles similar in mass to the top quark, not from an s-channel heavy resonance. It should, thus, be compared to the value $A_{t\bar{t}} \simeq 0.212 \pm 0.096$ for the reconstructed $M_{t\bar{t}} > 450$ GeV as actually measured within the CDF detector coverage [2].

We also check the total cross section of $t\bar{t}$ production at tree level. It is found to be within 10% of the value as calculated in the SM at the leading order. Assuming the same k-factor as in the case of the SM, we predict the total $t\bar{t}$ cross-section within theoretical uncertainties [27–30]. Deviation is seen when we check the cross-section as a function $M_{t\bar{t}}$. In Fig. 2 we have compared the tree level cross-section, as calculated in our model, to that in the SM. As shown in Fig. 2 and as reported in Refs. [40], the deviation is too small to give clear signal especially in early LHC data. Note that we cannot make direct comparison to the experimental data, since higher order corrections may change the shape of the curve besides the overall cross-section.

V. THE CDF EXCESS IN Wjj

The relevant interactions for these processes are:

$$\mathcal{L}_{\rm int} = -\lambda(\bar{c}_{\rm L}\Phi_3^0 u_{\rm R} + \bar{s}_{\rm L}\Phi_3^- u_{\rm R} + {\rm c.c.})\,. \qquad (11)$$

Wjj final states via an intermediate Φ then dominantly arise at the Tevatron from the processes $u\bar{s} \to W^+\Phi_3^0 \to W^+u\bar{c}$ and its charge-conjugated process.

For $\lambda = 1.4$, the Wjj signal cross-section is found to be 2.1 pb. We use `Madgraph v4.4.48` to generate the signal events, which are subsequently decayed, showered, and finally hadronized by `PYTHIA v6.4` [41]. We group the hadronic output of `PYTHIA` into "cells" of size $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$. We sum the four momentum of all particles in

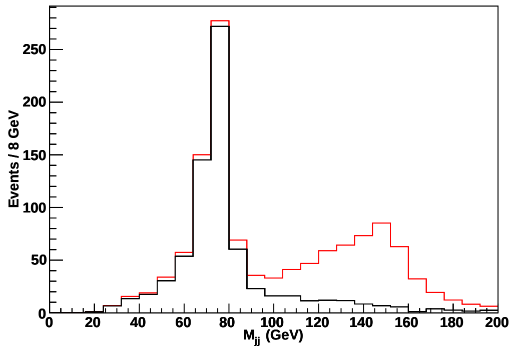

FIG. 3: The dijet invariant mass distribution. The red curve is due to the resonances in our model at 160 GeV. The black curve shows the same for SM diboson events. Note that the center of the peak is shifted to 150 GeV.

each cell and rescale the resulting three-momentum such as to make cells massless. Jet clustering is done using the CDF version of Run-II iterative cone algorithm with midpoint seeds, as implemented in `Fastjet` [42]. We do not perform any realistic detector simulation.

We demand exactly one isolated lepton with $p_{\mathrm{T}} > 20$ GeV and $|\eta| < 1.0$ plus missing transverse energy $\not{E}_{\mathrm{T}} > 25$ GeV. The event also contains exactly two jets, each with $E_{\mathrm{T}} > 30$ GeV and within $|\Delta\eta| < 2.5$ of each other. The dijet invariant mass distribution is plotted in Fig. 3. In order to get an idea of the size of the bump in our model, we have also plotted SM diboson events $(ZW^{\pm}, W^+W^-)$ that pass all our selection criteria. Note that the bump has all the features as seen in the data, including the size and the position of the peak at 145–150 GeV.

The CDF has used a Gaussian profile function of narrow width to fit the excess. On the contrary, a quick glance at the the dijet mass distribution in data suggests that data favors a broader peak. The dijet mass distribution in the range 170–220 GeV is characterized by upward fluctuations in each bin. Under modest smearing, we find that the position and the shape of the peak in our signal events remain relatively unaltered and slight excesses are generated in higher bins. A true comparison, however, can only be made after a decent detector simulation, which is beyond the scope of this paper.

One might expect that we should also see a similar excess in Zjj sample at or around 120–160 GeV. Interestingly, Zjj production via an intermediate Φ^0 dominantly arises from $u\bar{c} \to Z\Phi_3^0 \to Zu\bar{c}$. This process is, however, suppressed because of the c PDF, leading to a much smaller cross-section $\simeq 0.3$ pb, consistent with the absence of observation of such an excess around 160 GeV. For the $u\bar{s}$ initial state, the process must be $u\bar{s} \to Z\Phi_3^+ \to Zu\bar{s}$, which would lead to an excess in a different region (around 220 GeV). However, the cross-section for it is smaller ($\simeq 0.24$ pb) and the peak would appear even broader.

VI. CONCLUSIONS

The $t\bar{t}$ forward-backward asymmetry and the 3.2σ excess in the 120–160 GeV range of the dijet mass distribution in the Wjj sample at the Tevatron, if real, signify the existence of new physics at the electroweak scale. We constructed a simple, weakly-coupled renormalizable theory with one multiplet of scalar particles obeying the $\left(\prod_{i=1}^{3} \mathrm{U}(1)_{q_{\mathrm{L}i}} \times \mathrm{U}(1)_{u_{\mathrm{R}i}}\right) \times \mathrm{U}(3)_d \times \mathbb{Z}_3$ flavor symmetry, which ensures that the only source of flavor/CP violations is V_{CKM}. We showed that the model can explain the two anomalies in terms of a single mass and a single coupling constant, without conflicting with existing bounds.

The flavor symmetry of the model also makes definite predictions on the amount of forward-backward asymmetries in $c\bar{c}$ and $b\bar{b}$. The $c\bar{c}$ asymmetry arises predominantly from the $u\bar{u}$ initial state via the t-channel exchange of Φ_3^0, while the $b\bar{b}$ asymmetry is dominated by $c\bar{c}$ via Φ_1^- exchange. Therefore, the $c\bar{c}$ asymmetry is predicted to be comparable to the $t\bar{t}$ asymmetry, while the $b\bar{b}$ asymmetry is expected to be suppressed due to the smallness of the c parton distribution function.

As mentioned in Sec. V, we also have a Zjj production via $u\bar{s} \to Z\Phi_3^+ \to Zjj$, with smaller cross-section than Wjj. For the LHC with $\sqrt{s} = 7$ TeV, the production cross-sections of Wjj and Zjj are 47 pb and 12 pb respectively for the 160 GeV resonances, and are 20 pb and 8 pb respectively for the 220 GeV resonances.

Finally, note that we get a sizeable forward-backward asymmetry because of $\mathcal{O}(1)$ coupling of the top quark with scalars of masses comparable to top mass, and as a result, we expect to see larger t-$\bar{t}$ production cross section that in the SM at high values of the invariant mass of the t-$\bar{t}$ pairs. This would certainly be an interesting feature to observe at the LHC.

Acknowledgments

TSR acknowledges stimulating discussions with A. Martin. This project also benefited from conversations with S. Ellis, M. Schmaltz, and B. Tweedie. The work of AEN and TSR was supported, in part, by the US Department of Energy under contract numbers DE-FGO2-96ER40956.

[Note added] While this manuscript was in progress, Refs. [43–47] discussing the CDF Wjj excess (but not the $t\bar{t}$ asymmetry) appeared in arXiv.

[Note added 2] A week after the first version of our preprint appeared on arXiv, we have received Ref. [48], which claims that the Φ_- mass has to be heavier than 540 GeV to avoid too large a rate for $B^+ \to K^+\pi^0$, and that this would make our Wjj excess signal go away. These claims, however, are incorrect. First, even granting 540 GeV for the Φ^- mass, it is not true that the signal would disappear, since it comes from the production of

the scalar Φ^0 *directly from SM quarks,* whereas Ref. [48] misunderstands that the signal is from an s-channel Φ^- decaying to a Φ^0 and a W^-.

Second, the $B^+ \to K^+\pi^0$ (i.e. $\bar{b} \to d\bar{d}\bar{s}$) calculation of Ref. [48] is actually technically incorrect. The essential error is the misidentification of the first term of 4-fermion operators (7) as the standard QCD penguin operator O_6, where the latter is summed over all (active) quark flavors while the former is *not*. As a result, dangerous decays such as $\bar{b} \to d\bar{d}\bar{s}$ are only generated via renormalization group running and are small. A proper calculation analyzing $\overline{B}^0 \to \pi^+ K^-$ (i.e. $b \to u\bar{u}s$) was recently performed by Ref. [49], which found that the rate for this process is enhanced by two orders of magnitude in our model.

We propose two ways to avoid this constraint while keeping our signals intact. The first is simply to make Φ^- heavier by a factor of a few to suppress operators (7). This will enhance the T parameter (8) but it can be tuned to be small by adding, for example, an electroweak-triplet scalar with a nonzero vacuum expectation value which contributes negatively to T.

The second way, which would require no tuning, is to enlarge the $\mathbb{Z}_3$ symmetry of the flavor symmetry (1) to the full permutation group S_3. This amounts to the following. To keep track of S_3 more easily, we rename $\Phi_{1,2,3}$ as $\Phi_{32,13,21}$, respectively, and introduce $\Phi_{23,31,12}$ with the same gauge quantum numbers as $\Phi_{32,13,21}$ and the $\left(\prod_{i=1}^{3} \mathrm{U}(1)_{q_{\mathrm{L}i}} \times \mathrm{U}(1)_{u_{\mathrm{R}i}}\right)$ charges in the manner obvious from the notation (e.g., Φ_{23} has charges $+1$ and -1 under $\mathrm{U}(1)_{q_{\mathrm{L}2}}$ and $\mathrm{U}(1)_{u_{\mathrm{R}3}}$, respectively, which is the opposite of Φ_{32}). The six Φ fields form a sextet of S_3, and thus have a common mass m_Φ and a Φ^0-Φ^- mass splitting parameter ξ in the lagrangian (4). The Yukawa couplings in Eq. (4) has to be generalized to be S_3 invariant, i.e.,

$$\mathcal{L}_{\mathrm{Yukawa}} = \lambda \sum_{\text{all permutations}} \bar{q}_{\mathrm{L}1} \Phi_{12} u_{\mathrm{R}2} + \text{c.c.}\,. \tag{12}$$

Then, instead of the dangerous operator $(\bar{q}_{\mathrm{L}2} u_{\mathrm{R}})(\bar{u}_{\mathrm{R}} q_{\mathrm{L}2})$ (the first one of Eq. (6)), we generate

$$\begin{aligned} &(\bar{q}_{\mathrm{L}2} u_{\mathrm{R}})(\bar{u}_{\mathrm{R}} q_{\mathrm{L}2}) + (\bar{q}_{\mathrm{L}3} u_{\mathrm{R}})(\bar{u}_{\mathrm{R}} q_{\mathrm{L}3}) \\ &= -(\bar{q}_{\mathrm{L}1} u_{\mathrm{R}})(\bar{u}_{\mathrm{R}} q_{\mathrm{L}1}) + \sum_{i=1}^{3} (\bar{q}_{\mathrm{L}i} u_{\mathrm{R}})(\bar{u}_{\mathrm{R}} q_{\mathrm{L}i})\,, \end{aligned} \tag{13}$$

in the gauge basis. Going to the mass basis, the 2nd term will remain flavor-diagonal by the unitarity of V, hence not contributing to $b \to u\bar{u}s$, while the 1st term gives the b-dependent operator

$$V_{ub}^* V_{ui} (\bar{b}_{\mathrm{L}} u_{\mathrm{R}})(\bar{u}_{\mathrm{R}} d_{\mathrm{L}i})\,. \tag{14}$$

Remarkably, this is much smaller than the dangerous operator in Eq. (7), which is only suppressed by $V_{cb}^* V_{ci}$. Therefore, the $b \to u\bar{u}s$ transition rate in this new model is suppressed by the same CKM factors as the tree level standard model contribution, and is sufficiently small even with a fairly light Φ^-. The doubling of Φ will increase the dijet cross-section but still within the uncertainties discussed in section III, especially due to the fact that λ in this model is smaller for the same values of the $t\bar{t}$ asymmetry. A detailed study of this model is under investigation.

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